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STATE-DEPENDENT PAIR CORRELATIONS IN NUCLEI

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1. INTRODUCTION

Some years ago a translationally invariant configuration interaction (TICI) method was developed as a way to perform microscopic calculations for the ground-state properties of nuclear systems [1,2]; this approach was inspired by the coupled cluster (CC) method and the calculations performed for the ${}^4\text{He}$ nucleus were encouraging enough to pursue with this method. In this paper we want to discuss work which is still in progress [3], namely the extension of the lowest order version (TICI2) to heavier systems and the inclusion of state-dependent correlations.

In the SUB(2) approximation of the CC method, the wave function is obtained by acting with the exponentiated one- and two-body correlation operators, $\exp(S_1 + S_2)$, on an appropriate reference state which takes care of the symmetry properties of the system; but if we want to preserve the translational invariance these two correlation operators cannot be independent and must be combined to create a new correlation

operator $S^{(1,2)}$. If we consider the exponential of this new operator at first order and convert it to coordinate representation we have the TICI2 wave function:

$$|\Psi\rangle = \sum_{i<j} f(r_{ij})|\Psi_0\rangle, \quad (1)$$

where the uncorrelated state $|\Psi_0\rangle$ must be a ground state harmonic-oscillator shell-model function to avoid unwanted contributions of the center-of-mass motion. It is worth noting that because we restrict ourselves to the linear term only in the exponential form of the CC wave function, the TICI2 method is fully variational and the results are rigorous upper bounds to the ground-state energy.

To deal with realistic nucleon-nucleon interactions we must give an operatorial structure to our correlation functions which will be taken of the same form as the interaction. The manner in which state dependence must be taken into account may be rather involved in some methods, but in TICI2 the extension is straightforward because we have no commutation problems due to the nature of the wave function. We can introduce directly the state-dependence in our wave function by allowing the f in Eq. (1) to be state-dependent

$$f_{ij} = \sum_p f^p(r_{ij})\Theta_{ij}^p, \quad (2)$$

where Θ^p stands for an appropriate set of operators. We will consider here two different sets: the first one having a V4 structure ($1, P^\sigma, P^\tau$ and $P^\sigma P^\tau$, where P^σ and P^τ are the spin- and isospin-exchange operators), and the second one with a V6 form, that is V4 plus the tensor channels S_{ij} and $S_{ij}P^\tau$.

In Section 2 we briefly sketch our methodology and we turn to the results obtained for both the V4 and V6 cases in Section 3 and Section 4 respectively, comparing with other many-body methods and exact results where possible. Finally some concluding remarks will be presented in Section 5.

2. METHODOLOGY

In this section we explain how required matrix elements are computed, both in V4 and V6 form. We discuss only the potential matrix elements since they are the more complicated, including up to six different nucleons. The other matrix elements needed, those related to the kinetic energy operator and the norm, can be computed in a similar fashion.

First of all our reference state $|\Psi_0\rangle$, will be an independent-particle harmonic oscillator shell-model wave function. This uncorrelated reference state is a single Slater determinant for the case of spin-isospin saturated systems and we will represent it by the occupied harmonic oscillator states. It will be convenient to work in a Cartesian basis where the states are labeled by three harmonic oscillator quantum numbers (n_x, n_y, n_z); in this basis we have used non-spherical reference states for the ^8Be and ^{12}C nuclei, $|(0,0,0)^4(0,0,1)^4\rangle$ and $|(0,0,0)^4(1,0,0)^4(0,1,0)^4\rangle$ respectively, which has proven to be better than considering a spherical reference state [3].

Our goal now is to compute the expectation value of the potential energy operator, which is given by

$$\langle V \rangle = \langle \Psi_0 | \sum_{i < j} f^*(ij) \sum_{k < l} V(kl) \sum_{m < n} f(mn) | \Psi_0 \rangle, \quad (3)$$

where the indices in parentheses identify the particles the corresponding operator is acting on. Many of the terms in the previous expression are similar, so a simple way of dealing with them is to make a diagrammatic analysis [2] and get the topologically distinct diagrams to expand the matrix elements.

We have to perform a summation over all occupied single-particle states but since we are dealing with saturated systems this sum is reduced to the harmonic oscillator states times the respective trace of the spin and isospin operators, where we must take account of the permutations. Each permutation is split into the space, spin and isospin parts, $(ij) \equiv (i^s j^s)(i^\sigma j^\sigma)(i^\tau j^\tau)$, and the spin and isospin permutations can be represented by products of spin and isospin exchange operators, P_{ij}^σ and P_{ij}^τ , and incorporated to the trace over spin and isospin. The final matrix element related to a given diagram is

$$\begin{aligned} & \langle \Psi_0 | f^*(ij) V(kl) f(mn) | \Psi_0 \rangle \\ &= \sum_{pqr} \sum_{i_1^s \dots i_N^s=1}^{A/4} \sum_P \epsilon_P \text{Tr} (\Theta_{ij}^p \Theta_{kl}^q \Theta_{mn}^r P_{P i_1^s}^\sigma \dots P_{P i_N^s}^\sigma P_{P i_1^s}^\tau \dots P_{P i_N^s}^\tau) \\ & \times \langle \phi_{i_1^s}(1) \dots \phi_{i_N^s}(N) | f^{p*}(r_{ij}) V^q(r_{kl}) f^r(r_{mn}) | \phi_{P i_1^s}(1) \dots \phi_{P i_N^s}(N) \rangle. \end{aligned} \quad (4)$$

In this expression the ϕ 's are the spatial parts of the single particle wave functions and the Θ^k form the basis of operators in which both the correlations and the potential have been expanded.

The computation of the required traces is greatly simplified if we work in the direct product representation since in this representation all the needed matrices, exchange and spin matrices, have the same structure, with only one non-null element per row and column. In this manner the storage and matrix operations can be done very efficiently and, although the number of required traces is very high, the trace calculation may be done at once from the very beginning and stored for later use. However it is worth noting that the number of traces required in the V6 case is much higher than in the V4 case due to the complicate structure of the tensor operator.

The last step in the calculation of the matrix elements is related to the spatial integrals. There is here a basic difference between the V4 and V6 case; in the former the operatorial structure does not interfere with the spatial part of the calculation, but in the later a spatial contribution to the integrals appears. This is due to the spatial coordinates included in the tensor operator and we must consider consequently a higher number of integrals. In order to perform the integrations we expand the correlation functions in a set of Gaussians, $f^p(r) = \sum_k c_k^p \exp(-\beta_k r^2)$, a method which has proven to be an accurate representation of the correlation function f when including negative and positive values of β [2]. The integrals are computed (see [4]) by using a recurrence relation for all integrations apart from those involving the coordinates in the potential; the remaining integral involving the potential is then calculated numerically.

Now we can calculate the matrix elements of the Hamiltonian and the norm, and once evaluated these elements the expectation value of the energy, which depends on the set of coefficients $\{c_m^r\}$, can be expressed as $E = c_m^r \mathcal{H}_{mn}^r c_n^s / c_m^r \mathcal{N}_{mn}^r c_n^s$. Optimizing with respect to c_m^r we obtain the desired result for the energy, which is the lowest eigenvalue of a generalized eigenvalue problem.

It is worth commenting on the special ${}^4\text{He}$ case. Its wave function is fully space-symmetric and so the action on the spin-exchange operators is equal to the isospin-exchange one, with opposite sign (space symmetry makes $P_{ij}^\sigma P_{ij}^\tau = -1$). Thus in this case the basis for the correlations is overcomplete and to avoid problems when solving the generalized eigenvalue problem we discard the redundant terms, keeping only the ones without isospin-exchange operators.

3. V4 RESULTS

Having described our methodology in the previous section we turn now to look to the results obtained using the V4 structure for both the correlation and the nucleon-nucleon interaction. We have calculated the ground state energies of the spin and isospin saturated nuclei within the 0p-shell, that is ${}^4\text{He}$, ${}^8\text{Be}$, ${}^{12}\text{C}$ and ${}^{16}\text{O}$. The potentials used are the following: the Brink-Boeker B1 potential [5], the modified Afnan-Tang MS3 potential [6] and the Malfiet-Tjon MT I/III and MT V potentials [7]. In Table 1 we show for the different nuclei and interactions the state-dependent and -independent binding energies, together with the respective values of α (the inverse length-scale of the harmonic oscillator). The values of α are the ones that minimize the binding energy for the state-independent (SI) case, where the correlation function used is a purely scalar one. The optimal values found in the SI calculation were employed afterwards to perform the state-dependent (SD) calculation, so the gain in energy when going from SI to SD case may be somewhat underestimated.

We can see in Table 1 that in the ${}^4\text{He}$ case the value of the energy does not change when passing from state-independent to state-dependent correlations with the B1 and MT V interactions. This is not surprising if we consider the structure of these two potentials: the MT V potential is a pure Wigner force and the B1 interaction is a Wigner plus Majorana potential. So long as the ${}^4\text{He}$ uncorrelated reference state is a fully space-symmetric state, there is no coupling with the spin part of the state-dependent correlation (as mentioned in the previous section the ${}^4\text{He}$ correlation only presents Wigner and Bartlett terms), hence the energy does not change in these two cases. However, there is a contribution of this spin correlation to the heavier systems with these interactions since their uncorrelated reference states are no longer space-symmetric.

The general result of including state-dependent correlations (with the two exceptions mentioned above) is a significant lowering in the energy of the systems considered. The biggest influence of state-dependence appears in the MS3 potential as one would expect since this is the only one of the interactions considered having a complete V4 structure. For this potential the gain in energy ranges between 40% in ${}^8\text{Be}$ and 10% in ${}^4\text{He}$. The B1 interaction has a similar behavior with a somewhat smaller increase in the energy (15%-25%) than in the MS3 case, and the MT I/III interaction presents an improvement that is roughly a half of the obtained with B1.

Table 1

Binding energies (in MeV) of $0p$ -shell nuclei for the interactions considered, B1, MS3, MT I/III and MT V. The label SI corresponds to a purely scalar correlation function and the label SD to a correlation function of V4 type. α is the harmonic oscillator parameter (in fm $^{-1}$).

Interaction	Nucleus	α	E_{SI}	E_{SD}
B1	^4He	0.729	37.86	37.86
	^8Be	0.595	49.18	61.30
	^{12}C	0.595	84.91	103.93
	^{16}O	0.602	145.94	167.30
MS3	^4He	0.713	25.41	27.99
	^8Be	0.582	26.26	37.30
	^{12}C	0.588	46.22	62.99
	^{16}O	0.596	85.56	105.64
MT I/III	^4He	0.741	29.45	30.81
	^8Be	0.660	46.67	52.67
	^{12}C	0.698	99.05	109.04
	^{16}O	0.744	194.10	207.52
MT V	^4He	0.741	29.45	29.45
	^8Be	0.869	129.25	130.23
	^{12}C	0.997	425.99	429.44
	^{16}O	1.078	966.65	973.67

Lastly, the MT V potential is the less sensitive to the introduction of state-dependent correlations, changing less than 1% when going from SI to SD case. This was expected too, being the MT V interaction of Wigner type.

Considering the different nuclei we see that the pattern is almost the same for all the interactions. The gain in energy is bigger in ^8Be , then followed by ^{12}C , ^{16}O and ^4He in decreasing order. The exception to this rule is the MT V potential where the increase is almost constant, slightly better in ^{12}C than in ^8Be .

Table 2 presents a comparison between our results and those obtained with other various techniques for the Malfliet-Tjon MT V interaction. This potential is a purely scalar interaction and it produces a strong binding (it does not actually saturate nuclear matter). The techniques reported in Table 2 refer to variational Monte Carlo calculations (VMC), fixed-node Green's function Monte Carlo (GFMC) methods, diffusion Monte Carlo methods (DMC), the Fermi hypernetted chain method based on a Jastrow wave function (FHNC), the integro-differential equation approach (IDEA) and the stochastic variational method (SVM).

One can appreciate in Table 2 that our results are close to the Monte Carlo ones, even in the ^{16}O case. As mentioned above the MT V interaction gives a very strong

Table 2

A comparison of the binding energies (in MeV) of ${}^4\text{He}$ and ${}^{16}\text{O}$ using various techniques with the Malfliet-Tjon MT V potential. TICI2 is the method presented here, VMC, GFMC and DMC are variational, Green's function and diffusion Monte Carlo methods respectively, FHNC is the Fermi hypernetted chain method, IDEA is the integro-differential equation approach and SVM is the stochastic variational method.

	${}^4\text{He}$	${}^{16}\text{O}$
TICI2(SI)	29.45	966.65
TICI2(SD)	29.45	973.67
VMC		1024 $\pm$ 5 [8] 1103 $\pm$ 1 [9] 1138.5 $\pm$ 0.2 [10]
GFMC [8]	31.3 $\pm$ 0.2	1194 $\pm$ 20
DMC	31.32 $\pm$ 0.02 [11]	1189 $\pm$ 1 [10]
FHNC/0 [12]		987/1152
FHNC//0 [13]		1059/1055
IDEA [14]	30.7-31.2	1021-1027
SVM [15]	31.360	

binding and we expect the ${}^{16}\text{O}$ to be a high density system when considering this interaction and, since our method only considers pair correlations, one could expect a great difference with the essentially "exact" results. In despite of this considerations our state-dependent value, 973.67 MeV, amounts more than 80% of the very accurate diffusion Monte Carlo result [10], 1189 $\pm$ 1 MeV. One can also see that our results are close to those obtained with the methods based on a Jastrow correlation operator, product over all pairs of particles of a scalar (state-independent) pair function of the interparticle distance. The VMC results in Table 2 have been obtained with different forms of single-particle wave functions and correlation functions, hence the three different values showed. On the other hand the Fermi hypernetted chain approximations [12,13,17] represent resummed cluster expansions of the same quantity, in which some cluster terms are summed to all orders but neglecting or approximating some elementary diagrams.

The TICI2 results with the MT V interaction are quite reasonable but one would expect better results with the other interactions, which are not so extremely unrealistic as the MT V is. In Table 3 we compare the binding energies of ${}^{12}\text{C}$ and ${}^{16}\text{O}$ obtained using the Brink-Boeker B1 potential. This interaction has only Wigner and Majorana terms. Again all the other results in the table (save the IDEA ones) come from Jastrow wave functions with scalar correlations. The row FAHT-4 corresponds to a 4th order cluster expansion of the Jastrow energy functional. If we consider first the ${}^{12}\text{C}$ values we realize that our state-independent result is better than the other results, being around the VMC and FAHT values, but 15 MeV up the FHNC ones.

Table 3

A comparison of the binding energies (in MeV) using various techniques with the Brink-Boeker B1 potential. For a description of the methods see previous table. FAHT is a cluster expansion.

	^{12}C	^{16}O
TIC12(SI)	84.91	145.94
TIC12(SD)	103.93	167.30
VMC [16]	82.9 ± 0.2	150.9 ± 0.3
FHNC/0	68.9 [17]	161.7 [17]
		168.2/167.2 [12]
FHNC-1	63.7 [17]	151.4 [17]
		150.4/152.4 [12]
FAHT-4 [16]	82.3	148.6
IDEA [14]	79.8-80.1	164.4-165.2

When going from state-independent to the state-dependent case we find that our result is more than 20 MeV up all the other values. The results are very good but most of this difference is due to the different reference states used. Whereas we have chosen a non-spherical uncorrelated reference state for ^8Be and ^{12}C the other techniques have used a spherical uncorrelated reference state. In the ^{16}O case the reference states used in all the methods are spherically symmetric, so we do not expect our results to present a great gap respect the others like they do in the ^{12}C case. Our state-independent results for ^{16}O are close to the VMC and FAHT ones but they are now far below the FHNC//0 ones. However it is worth noting that when going to the next approximation order within FHNC it results in a loose of binding energy, so the FHNC-1 results are again near our state-independent numbers. State-dependent correlations rise again the binding energy about 20 MeV, being this state-dependent result at the same level than the best FHNC results.

4. V6 RESULTS

Let us now include tensor correlations in the wave function. We have performed calculations with the Argonne potentials AV14 [18] and AV18 [19], the Gogny-Pires-De Tourreil potential (GPDT) [20] and the De Tourreil-Sprung Super Soft Core potential (SSC) [21]. In all the cases we have used only the V6 part of the interactions, ignoring the other terms related to the spin-orbit interaction and other different operators. The introduction of the tensor correlations in our scheme increases a lot the computer time needed in the calculations, as stated in Section 2, and we have not considered for the time being the time-consuming ^{16}O case. In Table 4 we can see the results obtained for the different interactions mentioned above and the value of the harmonic oscillator parameter that minimizes the binding energy. Both AV14

Table 4

Binding energies (in MeV) of 0p-shell nuclei using V6 correlations with the V6 part of the interactions considered, AV14, AV18, GPDt and SSC. α is the harmonic oscillator parameter (in fm⁻¹).

Interaction	Nucleus	α	E
AV14	⁴ He	0.59	14.77
	⁸ Be	0.43	9.26
	¹² C	0.41	10.50
AV18	⁴ He	0.61	15.40
	⁸ Be	0.47	11.13
	¹² C	0.46	14.96
GPDt	⁴ He	0.70	27.36
	⁸ Be	0.58	40.79
	¹² C	0.61	73.37
SSC	⁴ He	0.68	24.12
	⁸ Be	0.54	25.98
	¹² C	0.54	39.64

and AV18 give a very low value (round 15 MeV) for the ⁴He energy and surprisingly the ⁸Be and ¹²C are less bounded. The results obtained with the GPDt and SSC potentials are closer to the expected ones.

In order to put some light in these results we have studied deeper the ⁴He system, where we can compare with some variational and Green's function Monte Carlo results. In addition to the potentials used before we have also considered the Reid soft-core potential[22], which presents a much harder core than the previous ones. We have also performed a new calculation in the ⁴He system using a wave function which is an hybrid between our familiar TICI2 wave function and Jastrow-like one. This hybrid has been constructed multiplying the TICI2 wave function by a purely scalar Jastrow correlation function, chosen to be of Gaussian form, $1 + a \exp(-br^2)$, which is the more appropriate to incorporate the Jastrow correlation to our previous scheme. With this wave function we keep the simple operatorial structure of TICI2 but adding the flexibility of the Jastrow correlations.

Our TICI2 and Jastrow-TICI2 (J-TICI2) results are compared in Table 5 with variational Monte Carlo results and with Green's function Monte Carlo results where possible. Note that the Reid interaction in the TICI2 scheme produces a poorly bound ⁴He nucleus. As mentioned before the Reid potential has a harder core than the Argonne interactions, so it seems the accuracy of TICI2 results is strongly related to the shape of the interaction.

Due to the additive ansatz for the pair correlations in TICI2 there are pairs of nucleons for which the interaction is not screened by the correlation. Consequently

Table 5

Comparison between ${}^4\text{He}$ energy values (in MeV) for several V6 interactions, AV14, AV18, GPDT, SSC and REID, obtained with different methodologies. TICI2 and J-TICI2 are our calculations, VMC and GFMC are variational and Green's function Monte Carlo methods.

	TICI2	J-TICI2	VMC	GFMC
AV14	14.77	20.37	23.24 ± 0.08 23.10 ± 0.08 [23] 22.75 ± 0.10 [23]	24.79 ± 0.20 [23]
AV18	15.40	21.08	24.80 ± 0.09	
GPDT	27.37	27.58	27.71 ± 0.06	
SSC	24.12	26.74	29.20 ± 0.12	
REID	5.67	22.70	27.82 ± 0.12 26.00 ± 0.08 [23]	28.30 ± 0.12 [23]

the harder the core the lesser the binding energy we will get with TICI2. A glance to Table 5 confirms that this is exactly the situation. Both Argonne potentials have a similar strong core and we can see that the difference between TICI2 numbers and Monte Carlo ones in these cases are similar too. The Reid potential has the harder core in this set of interactions, and TICI2 results differ enormously from VMC or GFMC results. In contrast, for the GPDT potential, which has the softer core, TICI2 provides a binding energy in nice agreement with VMC one.

As mentioned above these results are reflecting the fact that not all the interactions between nucleons are screened by the correspondent correlation factor. To illustrate this we have introduced Jastrow correlations in the TICI2 scheme and we expect now an improvement in the binding energy related to the strength of the core. Looking again at the table we can see that the gain is about 17 MeV for Reid interaction when going from TICI2 to J-TICI2, but only 0.2 MeV in GPDT case. The SSC potential gets 2.6 extra MeV and the Argonne potentials gain about 5.6 MeV. The TICI2 value for the GPDT potential is a good one, near the VMC result, and it changes very little when introducing Jastrow correlations. On the opposite the TICI2 value for the Reid interaction is 23 MeV down the GFMC result and when adding Jastrow correlations it gains 17 MeV, being still far from Monte Carlo results.

5. CONCLUDING REMARKS

We have reported here the extension of the TICI2 method in two different directions. First we have set out to deal with heavier systems than ${}^4\text{He}$, including within our scheme the saturated nuclei within the $0p$ -shell and developing the needed machinery. As a second goal we have incorporated state-dependent pair correlations in

this method in both V4 and V6 form, in order to deal with more realistic interactions. In Section 3 we have showed our results obtained when using V4 correlations and interactions and the comparisons made with other techniques. These results show the advantage of working with non-spherical uncorrelated reference states, as it can be seen when comparing the values of the binding energy for the ^{12}C nucleus. On the other hand, when dealing with spherical nuclei like ^4He or ^{16}O , our V4 results are competitive with those obtained with other methodologies like FHNC. It is worth noting too that our results are true upper bounds to the ground-state energies due to the fully variational character of TICI2 method. We can conclude that at V4 level this method is working well and it is a good starting point.

The V6 landscape is a very different one. First of all the purpose of taking account of tensor correlations is to deal with so-called realistic or semi-realistic interactions, which are much more complicated than the V4 ones. These potentials are fitted to deuteron properties and two-nucleon scattering data and usually have a relatively hard core. The results shown in Section 4 are not very encouraging: our method fails here and is not able to get energies near the ones obtained with Monte Carlo techniques. The analysis of the results in the ^4He case, with the inclusion of an hybrid calculation which incorporates Jastrow correlations, shows that the TICI2 method does not seem able to control the hard core contributions, at least in its present form. The values obtained when including some Jastrow correlations indicate that our method suffers here due to the lack of correlations in it. We need more than this additive structure of pair correlations to properly treat the contributions of a hard core.

For future investigations the first idea is to come back to the origins, the CC method that inspired us, and consider not TICI2 but TICC2, taking the exponential of the correlation exactly. But in this way the wave function has now much more terms and a TICI2-like calculation is no longer feasible for heavy systems, specially if we want to include state-dependent correlations. Another possible way is to work completely inside the coupled cluster framework solving the CC2 equations but losing the variational character of the solutions; besides, we know that when solving the CC2 equations in configuration space the convergence may be very slow. On the other hand we do not want to change our coordinate space wave function, which is easy to extend to state-dependent correlations, thus we must learn to solve the coupled cluster equations in coordinate representation. This is the way we are tentatively exploring hoping this allows us to treat accurately the pair correlations.

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REFERENCES

- [1] R.F. Bishop, M.F. Flynn, M.C. Boscá, E. Buendía, and R. Guardiola: Phys. Rev. C 42, (1990) 1341; Condensed Matter Theories, vol.5, ed. V.C. Aguilara-Navarro (Plenum Press, New York, 1990) p.253; J. Phys. G 16 (1990) L61.
- [2] R.F. Bishop, E. Buendía, M.F. Flynn, and R. Guardiola: Condensed Matter Theories, vol. 6, eds. S. Fantoni and S. Rosati (Plenum Press, New York, 1991) p. 405; J. Phys. G 17 (1991) 857, G 18 (1992) 1157, G 19 (1993) 1663.
- [3] R. Guardiola, P.I. Moliner, J. Navarro, R.F. Bishop, A. Puente, and N.R. Walet: Nucl. Phys. A 609 (1996) 218.
- [4] R. Guardiola, Nucl. Phys. A 328 (1979) 490.
- [5] D.M. Brink and E. Boeker, Nucl. Phys. A 91 (1967) 1.
- [6] R. Guardiola, A. Faessler, H. Müther and A. Polls, Nucl. Phys. A 371 (1981) 79.
- [7] R.A. Malfliet and J.A. Tjon, Nucl. Phys. A 127 (1969) 161.
- [8] J.G. Zabolitzky and M.H. Kalos, Nucl. Phys. A 356 (1981) 114;
U. Helmbrecht and J.G. Zabolitzky, Nucl. Phys. A 442 (1985) 109;
J.G. Zabolitzky, K.E. Schmidt and M.H. Kalos, Phys. Rev. C 25 (1982) 111.
- [9] E. Buendía and R. Guardiola, Condensed Matter Theories, vol. 8, eds. L. Blum and F.B. Malik (Plenum Press, New York, 1993) p. 301.
- [10] S.A. Chin and E. Krotscheck, Nucl. Phys. A 560 (1993) 151.
- [11] R.F. Bishop, E. Buendía, M.F. Flynn and R. Guardiola, J. Phys. G 18 (1992) L21.
- [12] G. Co', A. Fabrocini, S. Fantoni and I.E. Lagaris, Nucl. Phys. A 549 (1992) 439.
- [13] E. Krotscheck, Nucl. Phys. A 465 (1987) 461.
- [14] R. Brizzi, M. Fabre de la Ripelle and M. Lassaut, Nucl. Phys. A 596 (1996) 199.
- [15] K. Varga and Y. Suzuki, Phys. Rev. C 52 (1995) 2885; Phys. Rev. A 53 (1996) 1907.
- [16] M.C. Boscá, E. Buendía and R. Guardiola, Phys. Lett. B 198 (1987) 312; Condensed Matter Theories, vol. 3, eds. J. Arponen, R.F. Bishop and M. Manninen (Plenum Press, New York, 1987) p. 101.
- [17] F. Arias de Saavedra, G. Co', A. Fabrocini and S. Fantoni, Nucl. Phys. A 605 (1996) 359.
- [18] R.B. Wiringa, R.A. Smith and T.L. Ainsworth, Phys. Rev. C 29 (1984) 1207.
- [19] R.B. Wiringa, V.G.J. Stoks and R. Schiavilla, Phys. Rev. C 51 (1995) 38.
- [20] D. Gogny, P. Pires and R. de Tourreil, Phys. Lett. B 32 (1970) 591.
- [21] R. de Tourreil and D.W.L. Sprung, Nucl. Phys. A 201 (1973) 193.
- [22] R.V. Reid, Ann. Phys. 50 (1968) 411.
- [23] J. Carlson, Phys. Rev. C 36 (1987) 2026; C 38 (1988) 1879.